

Effective restoration of chiral symmetry in excited mesons.

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Abstract

A fast restoration of chiral symmetry in excited mesons is demonstrated. A minimal "realistic" chirally symmetric confining model is used, where the only interaction between quarks is the linear instantaneous Lorentz-vector confining potential. Chiral symmetry breaking is generated via the nonperturbative resummation of valence quarks self-energy loops and the meson bound states are obtained from the Bethe-Salpeter equation. The excited mesons fall into approximate chiral multiplets and lie on the approximately linear radial and angular Regge trajectories, though a significant deviation from the linearity of the angular trajectory is observed.

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There are certain phenomenological evidences that in highly excited hadrons, both in baryons [1, 2, 3] and mesons [4, 5] chiral and $U(1)_A$ symmetries are approximately restored, for a short overview see [6]. This "effective" restoration of chiral and $U(1)_A$ symmetries should not be confused with the chiral symmetry restoration at high temperatures and/or densities. What actually happens is that the excited hadrons gradually decouple from the quark condensates. Fundamentally it happens because in the high-lying hadrons the semiclassical regime is manifest and semiclassically quantum fluctuations of the quark fields are suppressed relative to the classical contributions which preserve both chiral and $U(1)_A$ symmetries [6, 7]. The microscopical reason is that in high-lying hadrons a typical momentum of valence quarks is large and hence they decouple from the quark condensate and consequently their Lorentz-scalar dynamical mass asymptotically vanishes [1, 3, 7, 10, 14]. Restoration of chiral symmetry requires a decoupling of states from the Goldstone bosons [3, 8, 9, 10] which is indeed observed phenomenologically since the coupling constant for $h^* \rightarrow h + \pi$ decreases fast higher in the spectrum.

At the moment there are two main paths to understand this phenomenon. In the first one one tries to connect the high-lying states to the short-range part of the two-point correlation function where the Operator Product Expansion is valid [2, 11]. However, the OPE is an asymptotic expansion. Then, while the correct spectrum of QCD must be consistent with the OPE, there is an infinite amount of incorrect spectra that can also be in agreement with the OPE. Hence the results within the present approach crucially depend on additional assumptions [12, 13].

In the second approach the authors try to understand this phenomenon within the microscopical models [7, 8, 10, 14]. There are also interesting attempts to formulate the problem on the lattice [16, 17], though extraction of the high-lying states on the lattice is a task of future.

In the absence of the controllable analytic solutions to QCD an insight into phenomenon can be achieved only through models. Clearly the model must be field-theoretical (in order to be able to exhibit the spontaneous breaking of chiral symmetry), chirally symmetric and contain confinement. In principle any possible gluonic interaction can contribute to chiral symmetry breaking and it is not known which specific interaction is the most important one in this respect. However, at the first stage it is reasonable to restrict oneself to the simplest possible model that contains all three key elements. Such a model is known, it

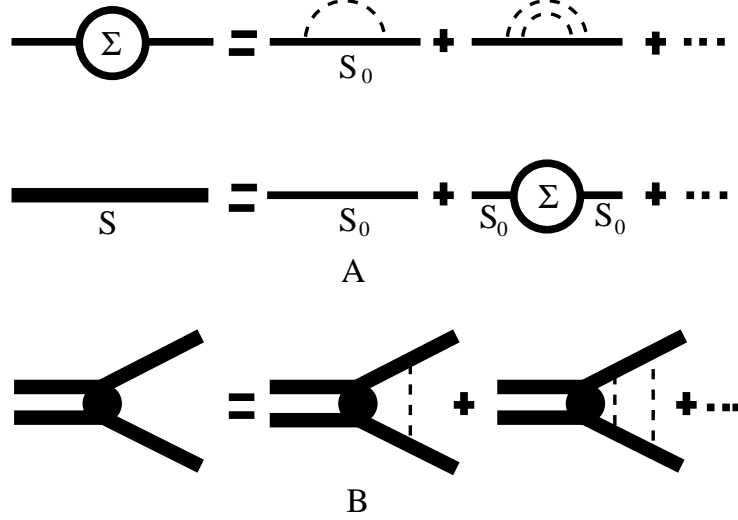


FIG. 1: Graphical representation of the Schwinger-Dyson (A) and Bethe-Salpeter equations (B) in the ladder approximation.

is a generalized Nambu and Jona-Lasinio model (GNJL) with the instantaneous Lorentz-vector confining kernel [18, 19, 20]. This model is similar in spirit to the large N_c 't Hooft model (QCD in 1+1 dimensions) [21]. In both models the only interaction between quarks is the instantaneous infinitely raising Lorentz-vector linear potential. Then chiral symmetry breaking is described by the standard summation of the valence quarks self-interaction loops in the rainbow approximation (the Schwinger-Dyson or gap equations), while mesons are obtained from the Bethe-Salpeter equation for the quark-antiquark bound states, see Fig. 1.

Conceptually the underlying physics is very clear in the 't Hooft model in the sense that once the proper gauge is chosen, the linear Lorentz-vector confining potential appears automatically as the Coulomb interaction in 1+1 dimensions. In 3+1 dimensions, once the Coulomb gauge is used for the gluonic field [22], an almost linearly raising confinement potential has been obtained [23, 24, 25].

An obvious advantage of the GNJL model is that it can be applied in 3+1 dimensions to systems of arbitrary spin. In 1+1 dimensions there is no spin, the rotational motion of quarks is impossible, and the states are characterized by the only quantum number, which is the radial quantum number. Then it is known that the spectrum represents an alternating sequence of positive and negative parity states and chiral multiplets never emerge. This happens because in 1+1 dimensions the valence quarks can perform only an oscillatory

motion. In 3+1 dimension, on the contrary, the quarks can rotate and hence can always be ultrarelativistic and chiral multiplets should emerge naturally [3].

Restoration of chiral symmetry in excited heavy-light mesons has been studied with the quadratic confining potential [14] and it was also mentioned in a model with the instantaneous potential of a more complicated form [15]. Here we report our results for excited light-light mesons with the linear potential. To our best knowledge this is the first explicit demonstration of the restoration of the chiral symmetry within the solvable "realistic" field-theoretical model. The results for the lowest mesons within the similar model have been previously reported in ref. [26].

The GNJL model is described by the Hamiltonian [19]

$$\begin{aligned}\hat{H} = & \int d^3x \bar{\psi}(\vec{x}, t) \left(-i\vec{\gamma} \cdot \vec{\nabla} + m \right) \psi(\vec{x}, t) \\ & + \frac{1}{2} \int d^3x d^3y J_\mu^a(\vec{x}, t) K_{\mu\nu}^{ab}(\vec{x} - \vec{y}) J_\nu^b(\vec{y}, t),\end{aligned}\quad (1)$$

with the quark current-current ($J_\mu^a(\vec{x}, t) = \bar{\psi}(\vec{x}, t) \gamma_\mu \frac{\lambda^a}{2} \psi(\vec{x}, t)$) interaction parametrised by an instantaneous confining kernel $K_{\mu\nu}^{ab}(\vec{x} - \vec{y})$ of a generic form. In this paper, we use the linear confining potential,

$$K_{\mu\nu}^{ab}(\vec{x} - \vec{y}) = g_{\mu 0} g_{\nu 0} \delta^{ab} V_0(|\vec{x} - \vec{y}|), \quad (2)$$

and absorb the color Casimir factor into string tension, $\frac{\lambda^a \lambda^a}{4} V_0(r) = \sigma r$.

The Schwinger-Dyson equation for the self-energy operator $\Sigma(\vec{p}) = [A_p - m] + (\vec{\gamma} \hat{\vec{p}}) [B_p - p]$ is

$$i\Sigma(\vec{p}) = \int \frac{d^4k}{(2\pi)^4} V(\vec{p} - \vec{k}) \gamma_0 \frac{1}{S_0^{-1}(k_0, \vec{k}) - \Sigma(\vec{k})} \gamma_0, \quad (3)$$

where

$$V(\vec{p}) = - \int d^3x e^{i\vec{p}\vec{x}} \sigma |\vec{x}| = \frac{8\pi\sigma}{p^4}, \quad (4)$$

so that the dressed Dirac operator becomes

$$D(p_0, \vec{p}) = iS^{-1}(p_0, \vec{p}) = \gamma_0 p_0 - (\vec{\gamma} \hat{\vec{p}}) B_p - A_p, \quad (5)$$

where, due to the instantaneous nature of the interaction the time-component of the Dirac operator is not dressed. The Lorentz-scalar dynamical mass A_p as well as the Lorentz-vector spatial part B_p contain both the classical and quantum contributions, the latter coming from

loops [7]:

$$A_p = m + \frac{1}{2} \int \frac{d^3 k}{(2\pi)^3} V(\vec{p} - \vec{k}) \sin \varphi_k, \quad (6)$$

$$B_p = p + \frac{1}{2} \int \frac{d^3 k}{(2\pi)^3} (\hat{\vec{p}} \cdot \hat{\vec{k}}) V(\vec{p} - \vec{k}) \cos \varphi_k, \quad (7)$$

where $\tan \varphi_p = \frac{A_p}{B_p}$.

Solution of the Schwinger-Dyson equation (3) with the linear potential is well-known, see e.g. [27], and the mass-gap equation has a nontrivial solution which breaks chiral symmetry, by generating a nontrivial dynamical mass function A_p . This dynamical mass is a very fast decreasing function at larger momenta. Then the quark condensate is given as

$$\langle \bar{q}q \rangle = -\frac{N_C}{\pi^2} \int_0^\infty dp p^2 \sin \varphi_p. \quad (8)$$

The homogeneous Bethe-Salpeter equation for the quark-antiquark bound state with mass M in the rest frame, i.e. with the four momentum $P^\mu = (M, \vec{P} = 0)$, is

$$\begin{aligned} \chi(\vec{p}, M) = & -i \int \frac{d^4 k}{(2\pi)^4} V(\vec{p} - \vec{k}) \gamma_0 S(k_0 + M/2, \vec{k}) \\ & \times \chi(\vec{k}, M) S(k_0 - M/2, \vec{k}) \gamma_0, \end{aligned} \quad (9)$$

where $\chi(\vec{p}, M)$ is the mesonic Salpeter amplitude in the rest frame. Eq. (9) is written in the ladder approximation for the vertex which is consistent with the rainbow approximation for the quark mass operator and which is well justified in the large- N_C limit.

The Salpeter amplitude can be decomposed into two components for mesons with $J^{PC} = (2n)^{-+}, (2n+1)^{+-}, (2n+1)^{++}, (2n+2)^{--}$, or 0^{++} and into four components for mesons with $J^{PC} = (2n+1)^{--}$ or $(2n+2)^{++}$, respectively. Here J is the spin, P the parity and C the charge conjugation parity of the meson and $n \in \mathbb{N}_0$. In that way the Bethe-Salpeter equation becomes a system of coupled integral equations for the components which we solve by expanding them into a finite number of properly chosen basis functions. This leads to a matrix eigenvalue problem which can be solved by standard linear algebra methods. We vary the meson mass until one of the eigenstates is equal to one.

In the gap as well as in the Bethe-Salpeter equations the infrared divergences are removed by introducing a finite "mass" into a confining potential, which is a standard trick. Then in the infrared limit ("mass" goes to 0) the quark propagator consists of a finite and diverged parts, while the mesons masses are finite. Recently it was demonstrated that also the masses

of quark-quark subsystems in the color-antitriplet state go to infinity in this limit and hence are removed from the physical spectrum [28]. The results presented in the following were obtained by calculating the infrared limit numerically, i.e. the quoted meson masses were extrapolated to the infrared limit from a few points with a very small but finite mass of the infrared regulator. It turned out that in this region of the small mass of the infrared regulator the squares of the meson masses depend almost linearly on the mass of the infrared regulator making the extrapolation reliable. The presented results are accurate within the quoted digits at least for states with small J and for states with higher J but small n . At larger J for larger n numerical errors accumulate in the second digit after comma.

By definition an effective chiral symmetry restoration means that (i) the states fall into approximate multiplets of $SU(2)_L \times SU(2)_R$ and the splittings within the multiplets ($\Delta M = M_+ - M_-$) vanish at $n \rightarrow \infty$ and/or $J \rightarrow \infty$; (ii) the splitting within the multiplet is much smaller than between the two subsequent multiplets [4, 5, 6].

The condition (i) is very restrictive, because the structure of the chiral multiplets for the $J = 0$ and $J > 0$ mesons is very different [4, 5]. For the $J > 0$ mesons chiral symmetry requires a *doubling* of states with some quantum numbers in contrast to the $J = 0$ states. Given the complete set of standard quantum numbers I, J^{PC} , the multiplets of $SU(2)_L \times SU(2)_R$ are

$$\mathbf{J} = \mathbf{0}$$

$$\begin{aligned} (1/2, 1/2)_a &: 1, 0^{-+} \longleftrightarrow 0, 0^{++} \\ (1/2, 1/2)_b &: 1, 0^{++} \longleftrightarrow 0, 0^{-+}, \end{aligned} \tag{10}$$

$$\mathbf{J} = 2\mathbf{k}, \quad \mathbf{k}=1,2,\dots$$

$$\begin{aligned} (0, 0) &: 0, J^{--} \longleftrightarrow 0, J^{++} \\ (1/2, 1/2)_a &: 1, J^{-+} \longleftrightarrow 0, J^{++} \\ (1/2, 1/2)_b &: 1, J^{++} \longleftrightarrow 0, J^{-+} \\ (0, 1) \oplus (1, 0) &: 1, J^{++} \longleftrightarrow 1, J^{--} \end{aligned} \tag{11}$$

$$\mathbf{J} = 2\mathbf{k}-1, \quad \mathbf{k}=1,2,\dots$$

$$\begin{aligned}
(0, 0) &: 0, J^{++} \longleftrightarrow 0, J^{--} \\
(1/2, 1/2)_a &: 1, J^{+-} \longleftrightarrow 0, J^{--} \\
(1/2, 1/2)_b &: 1, J^{--} \longleftrightarrow 0, J^{+-} \\
(0, 1) \oplus (1, 0) &: 1, J^{--} \longleftrightarrow 1, J^{++}
\end{aligned} \tag{12}$$

Note that within the present model the axial anomaly is absent. Even so there are no exact $U(1)_A$ multiplets, because this symmetry is broken not only by the anomaly, but also by the chiral condensate of the vacuum. Then the mechanism of the $U(1)_A$ symmetry breaking and restoration is exactly the same as of $SU(2)_L \times SU(2)_R$. Hence the effective restoration of $SU(2)_L \times SU(2)_R$ would automatically imply restoration of $U(1)_A$ and of $U(2)_L \times U(2)_R$ and vice versa. An effective restoration of the $U(1)_A$ symmetry would mean an approximate degeneracy of the opposite spatial parity states with the same isospin from the distinct $(1/2, 1/2)_a$ and $(1/2, 1/2)_b$ multiplets of $SU(2)_L \times SU(2)_R$.

Note that within the present model there are no vacuum fermion loops. Then since the interaction between quarks is flavor-blind the states with the same J^{PC} but different isospins from the distinct multiplets $(1/2, 1/2)_a$ and $(1/2, 1/2)_b$ as well as the states with the same J^{PC} but different isospins from $(0, 0)$ and $(0, 1) \oplus (1, 0)$ representations are exactly degenerate. Hence it is enough to show a complete set of the isovector (or isoscalar) states.

In Table 1 we present our results for the spectrum for the two-flavor (u and d) mesons in the chiral limit. Clearly the model should not be taken seriously for the low-lying states where other gluonic interactions as well as the $1/N_c$ corrections should be important. The purpose of the study is, however, to demonstrate that a solvable field-theoretical model in 3+1 dimensions does exhibit the effective restoration of the chiral symmetry at large radial excitations n and large spins. The excited mesons fall into approximate chiral and $U(1)_A$ multiplets and all conditions of the effective symmetry restorations are satisfied. We observe a very fast restoration of both $SU(2)_L \times SU(2)_R$ and $U(1)_A$ symmetries with increasing J and essentially more slow restoration with increasing of n .

When the chiral symmetry breaking Lorentz-scalar dynamical mass of quarks is zero, then there are independent Bethe-Salpeter amplitudes just according to the chiral representations (10)-(12). A finite dynamical mass plays a role of the off-diagonal matrix element and mixes the otherwise independent chiral Bethe-Salpeter amplitudes for the states $1^{--}, 2^{++}, 3^{--}, \dots$

TABLE I: Masses of isovector mesons in units of $\sqrt{\sigma}$.

| chiral multiplet | J^{PC} | radial excitation n | | | | | | | |
|------------------------|----------|-----------------------|------|------|------|------|------|------|--|
| | | 0 | 1 | 2 | 3 | 4 | 5 | 6 | |
| $(1/2, 1/2)_a$ | 0^{-+} | 0.00 | 2.93 | 4.35 | 5.49 | 6.46 | 7.31 | 8.09 | |
| $(1/2, 1/2)_b$ | 0^{++} | 1.49 | 3.38 | 4.72 | 5.80 | 6.74 | 7.57 | 8.33 | |
| $(1/2, 1/2)_a$ | 1^{+-} | 2.68 | 4.03 | 5.15 | 6.14 | 7.01 | 7.80 | 8.53 | |
| $(1/2, 1/2)_b$ | 1^{--} | 2.78 | 4.18 | 5.32 | 6.30 | 7.17 | 7.96 | 8.68 | |
| $(0, 1) \oplus (1, 0)$ | 1^{--} | 1.55 | 3.28 | 4.56 | 5.64 | 6.57 | 7.40 | 8.16 | |
| $(0, 1) \oplus (1, 0)$ | 1^{++} | 2.20 | 3.73 | 4.95 | 5.98 | 6.88 | 7.69 | 8.43 | |
| $(1/2, 1/2)_a$ | 2^{-+} | 3.89 | 4.98 | 5.94 | 6.80 | 7.59 | 8.31 | 8.99 | |
| $(1/2, 1/2)_b$ | 2^{++} | 3.91 | 5.02 | 6.00 | 6.88 | 7.67 | 8.41 | 9.09 | |
| $(0, 1) \oplus (1, 0)$ | 2^{++} | 3.60 | 4.67 | 5.64 | 6.51 | 7.31 | 8.06 | 8.75 | |
| $(0, 1) \oplus (1, 0)$ | 2^{--} | 3.67 | 4.80 | 5.80 | 6.68 | 7.49 | 8.23 | 8.91 | |
| $(1/2, 1/2)_a$ | 3^{+-} | 4.82 | 5.77 | 6.62 | 7.41 | 8.13 | 8.81 | 9.45 | |
| $(1/2, 1/2)_b$ | 3^{--} | 4.82 | 5.78 | 6.65 | 7.44 | 8.17 | 8.86 | 9.50 | |
| $(0, 1) \oplus (1, 0)$ | 3^{--} | 4.68 | 5.63 | 6.48 | 7.26 | 7.99 | 8.67 | 9.30 | |
| $(0, 1) \oplus (1, 0)$ | 3^{++} | 4.69 | 5.66 | 6.53 | 7.32 | 8.06 | 8.75 | 9.39 | |
| $(1/2, 1/2)_a$ | 4^{-+} | 5.59 | 6.45 | 7.23 | 7.96 | 8.64 | 9.28 | 9.89 | |
| $(1/2, 1/2)_b$ | 4^{++} | 5.59 | 6.45 | 7.24 | 7.97 | 8.66 | 9.30 | 9.92 | |
| $(0, 1) \oplus (1, 0)$ | 4^{++} | 5.51 | 6.36 | 7.15 | 7.88 | 8.56 | 9.20 | 9.80 | |
| $(0, 1) \oplus (1, 0)$ | 4^{--} | 5.51 | 6.37 | 7.16 | 7.90 | 8.58 | 9.23 | 9.84 | |
| $(1/2, 1/2)_a$ | 5^{+-} | 6.27 | 7.05 | 7.78 | 8.47 | 9.11 | 9.72 | 10.3 | |
| $(1/2, 1/2)_b$ | 5^{--} | 6.27 | 7.06 | 7.79 | 8.47 | 9.12 | 9.73 | 10.3 | |
| $(0, 1) \oplus (1, 0)$ | 5^{--} | 6.21 | 7.00 | 7.73 | 8.41 | 9.06 | 9.67 | 10.3 | |
| $(0, 1) \oplus (1, 0)$ | 5^{++} | 6.21 | 7.00 | 7.73 | 8.42 | 9.07 | 9.68 | 10.3 | |
| $(1/2, 1/2)_a$ | 6^{-+} | 6.88 | 7.61 | 8.29 | 8.94 | 9.55 | 10.1 | 10.7 | |
| $(1/2, 1/2)_b$ | 6^{++} | 6.88 | 7.61 | 8.29 | 8.94 | 9.56 | 10.1 | 10.7 | |
| $(0, 1) \oplus (1, 0)$ | 6^{++} | 6.83 | 7.57 | 8.25 | 8.90 | 9.51 | 10.1 | 10.7 | |
| $(0, 1) \oplus (1, 0)$ | 6^{--} | 6.83 | 7.57 | 8.26 | 8.90 | 9.52 | 10.1 | 10.7 | |

A key feature of this dynamical mass is that it is strongly momentum-dependent and vanishes very fast once the momentum is increased. When one increases excitation energy of a hadron, one also increases a typical momentum of valence quarks. Consequently, the chiral symmetry violating dynamical mass of quarks becomes small. Hence the mixing of the independent chiral Bethe-Salpeter amplitudes becomes small. A given state in the table is then assigned to the chiral representation according to the chiral Bethe-Salpeter amplitude that dominates in the given state.

In Fig. 2 the rates of the symmetry restoration against the radial quantum number n and spin J are shown. It is seen that with the fixed J the splitting within the multiplets ΔM decreases asymptotically as $1/\sqrt{n}$, dictated by the asymptotic linearity of the radial Regge trajectories. This property is consistent with the dominance of the free quark loop logarithm at short distances.

In Fig. 3 the angular and radial Regge trajectories are shown. Both kinds of trajectories exhibit deviations from the linear behavior. This fact is obviously related to the chiral symmetry breaking effects for lower mesons. Note, that the chiral symmetry requires a doubling of some of the radial and angular Regge trajectories for $J = 1, 2, \dots$. This is a highly nontrivial prediction of chiral symmetry. For example, some of the rho-mesons lie on the trajectory that is characterized by the chiral index $(0,1)+(1,0)$, while the other fit the trajectory with the chiral index $(1/2, 1/2)_b$. The intercepts of the asymptotic angular Regge trajectories for mesons in the given and different chiral representations coincide. Hence the asymptotic rate of the symmetry restoration with J is *faster* than $1/\sqrt{J}$.

The numerical result for the quark condensate is $\langle \bar{q}q \rangle = (-0.231\sqrt{\sigma})^3$, which agrees with the previous studies within the same model. If we fix the string tension from the phenomenological angular Regge trajectories, then $\sqrt{\sigma} \approx 300 - 400$ MeV and hence the quark condensate is between $(-70 \text{ MeV})^3$ and $(-90 \text{ MeV})^3$ which obviously underestimates the phenomenological value. Probably this indicates that other gluonic interactions could also contribute to chiral symmetry breaking. Notice that the string tension in Coulomb gauge can be larger than the asymptotic one. Lattice results suggest a value about twice the value obtained here [29, 30]. This would increase the value for the condensate but on the other hand lead to an unrealistically small pion decay constant [31, 32].

In the limit $n \rightarrow \infty$ and/or $J \rightarrow \infty$ one observes a complete degeneracy of all multiplets,

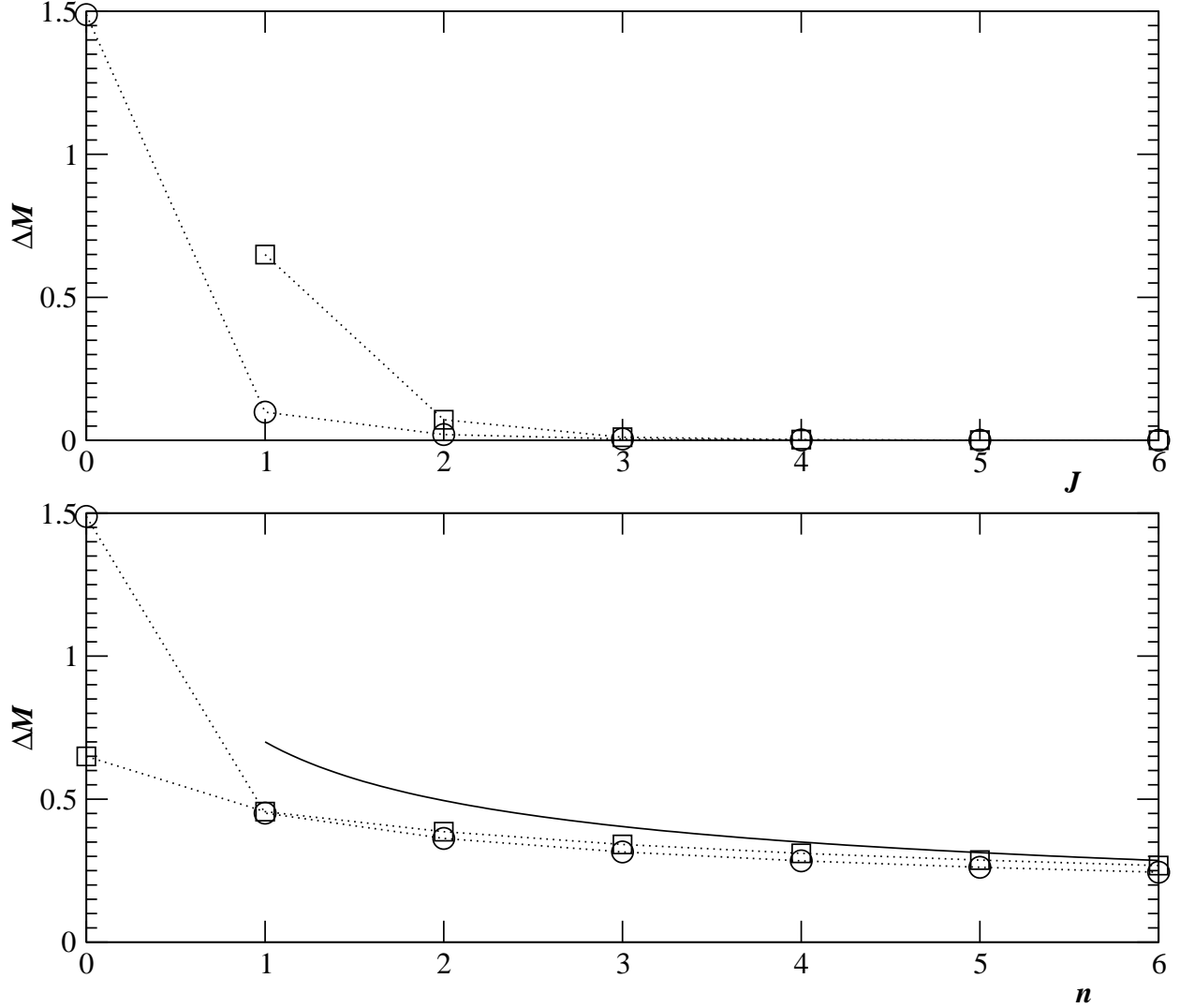


FIG. 2: Mass splittings in units of $\sqrt{\sigma}$ for isovector mesons of the chiral multiplets $(1/2, 1/2)_a$ and $(1/2, 1/2)_b$ (circles) and within the multiplet $(0, 1) \oplus (1, 0)$ (squares) against J for $n = 0$ (top) and against n for $J = 0$ and $J = 1$, respectively (bottom). The full line in the bottom plot is $0.7\sqrt{\sigma/n}$.

which means that the states fall into

$$[(0, 1/2) \oplus (1/2, 0)] \times [(0, 1/2) \oplus (1/2, 0)]$$

representation that combines all possible chiral representations for the systems of two massless quarks [5]. This means that in this limit the loop effects disappear completely and the system becomes classical [6, 7].

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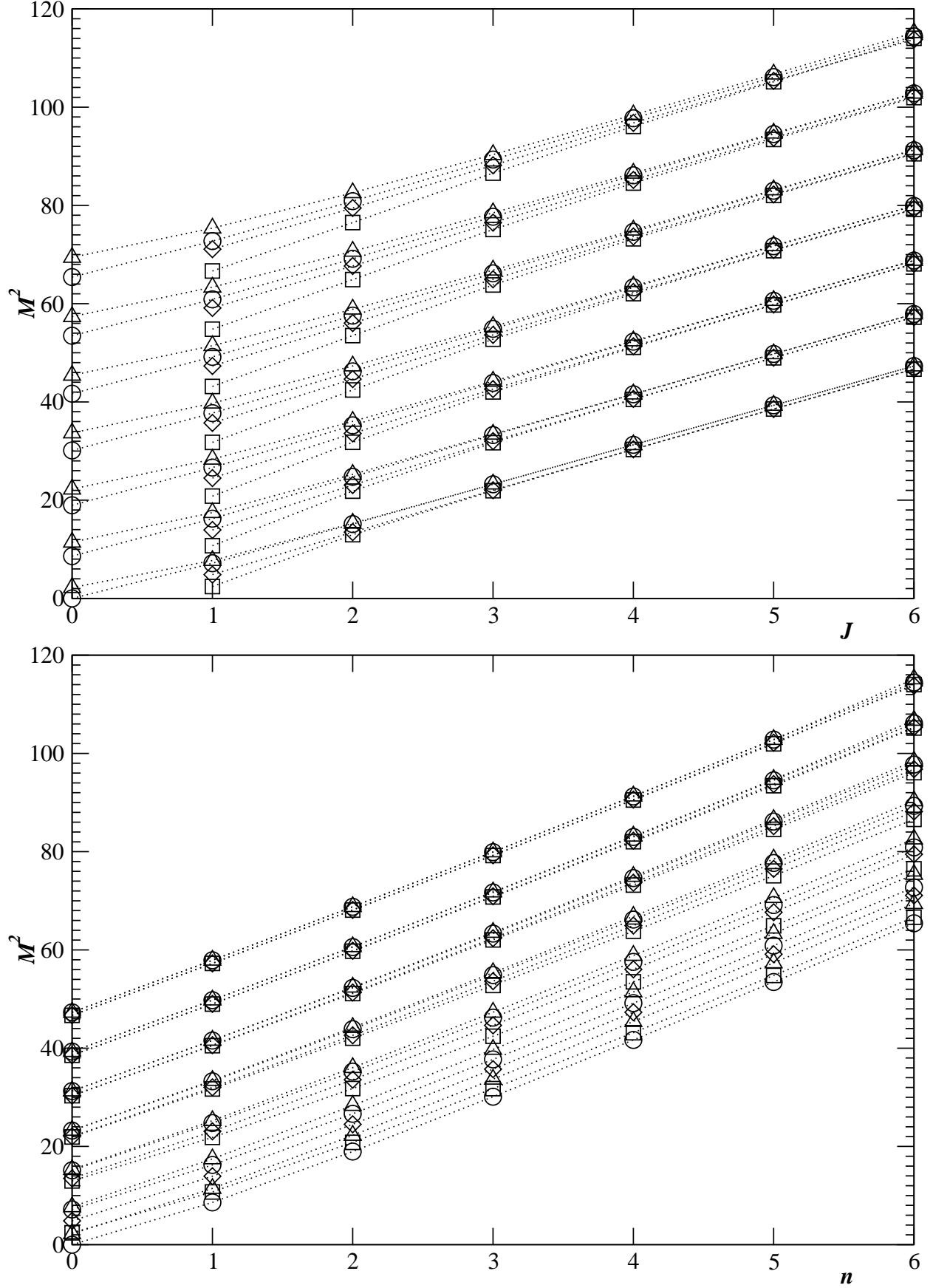


FIG. 3: Angular (top) and radial (bottom) Regge trajectories for isovector mesons with M^2 in units of σ . Mesons of the chiral multiplet $(1/2, 1/2)_a$ are indicated by circles, of $(1/2, 1/2)_b$ by

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- [1] L. Ya. Glozman, Phys. Lett. B **475** (2000) 329.
 - [2] T. D. Cohen and L. Ya. Glozman, Phys. Rev. D **65** (2002) 016006; Int. J. Mod. Phys. A **17** (2002) 1327.
 - [3] L. Ya. Glozman, Phys. Lett. B **541** (2002) 115.
 - [4] L. Ya. Glozman, Phys. Lett. B **539** (2002) 257.
 - [5] L. Ya. Glozman, Phys. Lett. B **587** (2004) 69.
 - [6] L. Ya. Glozman, Int. J. Mod. Phys. A. **21** (2006) 475.
 - [7] L. Ya. Glozman, A. V. Nefediev, J.E.F.T. Ribeiro, Phys. Rev. D **72** (2005) 094002.
 - [8] T. D. Cohen and L. Ya. Glozman, Mod. Phys. Lett. **A21** (2006) 1939.
 - [9] R. L. Jaffe, D. Pirjol, A. Scardicchio, hep-ph/0602010; Phys. Rev. D **74** (2006) 057901.
 - [10] L. Ya. Glozman, A. V. Nefediev, Phys. Rev. D **73** (2006) 074018.
 - [11] S. Beane, Phys. Rev. D **64** (2001) 116010; S.S. Afonin et al, J. High. Energy. Phys. **04** (2004) 039; E. Swanson, Phys. Lett. B **582** (2004) 167.
 - [12] M. Shifman, in: "Quark-Hadron Duality and the Transition to pQCD", Eds. A. Fantoni, S. Liuti, and O. Rondon (World Scientific, Singapore, 2006), pp. 171-192 (hep-ph/0507246).
 - [13] O. Cata, M. Golterman, and S. Peris, Phys. Rev. D **74** (2006) 016001.
 - [14] Yu. S. Kalashnikova, A. V. Nefediev, J.E.F.T. Ribeiro, Phys. Rev. D **72** (2005) 034020.
 - [15] N. Ligterink and E. S. Swanson, Phys. Rev. C **69** (2004) 025204.
 - [16] T. DeGrand, Phys. Rev. D **64** (2004) 074024.
 - [17] T. D. Cohen, Nucl. Phys. **A775** (2006) 89.
 - [18] J. R. Finger and J. E. Mandula, Nucl. Phys. B **199** (1982) 168.
 - [19] A. Le Yaouanc, L. Oliver, O. Pene, and J. C. Raynal, Phys. Rev. D **29** (1984) 1233; **31** (1985) 137.
 - [20] P. Bicudo and J. E. Ribeiro, Phys. Rev. D **42** (1990) 1611; **42** (1990) 1625.
 - [21] G. 't Hooft, Nucl. Phys. B **75** (1974) 461.
 - [22] N. H. Christ and T. D. Lee, Phys. Rev. D **22** (1980) 939.
 - [23] D. Zwanziger, Phys. Rev. D **70** (2004) 094034.
 - [24] A. P. Szczepaniak and E. Swanson, Phys. Rev. D **65** (2002) 025012.

- [25] C. Feuchter and H. Reinhardt, Phys. Rev. D **70** (2004) 105021; H. Reinhardt and C. Feuchter, Phys. Rev. D **71** (2005) 105002.
- [26] F. J. Llanes-Estrada and S. R. Cotanch, Phys. Rev. Lett., **84** (2000) 1102.
- [27] P. J. A. Bicudo and A. V. Nefediev, Phys. Rev. D **68** (2003) 065021.
- [28] R. Alkofer, M. Klokner, A. Krassnigg, and R. F. Wagenbrunn, Phys. Rev. Lett. **96** (2006) 022001.
- [29] J. Greensite, Š. Olejník, and D. Zwanziger, Phys. Rev. D **69** (2004) 074506.
- [30] A. Nakamura and T. Saito, Prog. Theor. Phys. **115** (2006) 189.
- [31] S. L. Adler and A. C. Davis, Nucl. Phys. B **244** (1984) 469.
- [32] R. Alkofer and P. A. Amundsen, Nucl. Phys. B **306** (1988) 305.